

# Associated Photons and New Physics Signals at Linear Colliders

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## Abstract

We study signals for beyond standard model physics and consider the virtues of single photon signals or associated photons in the final states in identifying different scenarios of new physics models in a very efficient and novel way.

*Keywords:* single photons; linear collider; resonances; physics beyond standard model.

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## 1 Introduction

Human scientific knowledge has come a long way in its quest for solutions to the unanswered questions and mysteries of nature, starting from the discovery of molecules and atoms to the present day knowledge of the smallest constituents of matter. And what we know today about the most fundamental building blocks of matter and the nature of forces governing their interaction is at best explained by the *Standard Model* (SM) of particle physics. Despite the remarkable successes, it is unlikely that the Standard Model is actually a complete theory of the fundamental laws of nature. Although the standard model is a mathematically consistent renormalizable field theory whose predictions have matched and withstood experimental tests down to at least  $10^{-16}$ cm, with an exception of the Higgs sector, it still leaves us with a lot many unresolved theoretical issues, like for instance the origin of mass, CP violation, number of fermion flavours, the hierarchy problems, etc.

Experimental hints for neutrino masses [1] already provide us with the necessity to consider physics beyond the SM. Another important issue with the SM is the stabilization of the electroweak scale and the

origin of mass. SM requires an elementary scalar in the theory which is responsible for the mass of all the particles in SM through the Higgs mechanism [2]. However, the physical state associated with this scalar, the Higgs boson, is yet to be experimentally observed. Electroweak precision data require the mass to be less than  $\sim$  few 100 GeV. Further more the Higgs mass squared receives large quantum corrections which drive its mass to the cut-off of the theory. One needs physics beyond the SM to stabilize the Higgs mass which is related to the scale of electroweak symmetry breaking in SM, in order to get the natural scale. Also the very fact that gravity is not fundamentally unified with the other interactions in the SM and there is no way to generate a quantum theory for gravity within the SM leads us to explore new theoretical ideas that extend beyond the SM and try to address physics issues concerning the SM. The nice feature of many of the proposed models is that they are predictive and should assert themselves at the TeV scale.

The world is sitting at this energy frontier with the running of Tevatron at Fermilab and the advent of the Large Hadron Collider (LHC) at CERN which would help us probe this disillusioned energy scale for any signal of new physics beyond the standard model and the origin of electroweak symmetry breaking (EWSB) with its associated mechanism that endows masses to the elementary particles. With a creative stream of tentative answers steadily flowing in and a distinct trait that seems common to all proposals, that some kind of new physics phenomenology must exist at the scale of TeV, pushing the SM as some form of an effective theory in the low energy limit we may be sitting at the cross-roads of discovery. It is expected that candidate theories like supersymmetry, technicolour, little Higgs and models of extradimensions, to name a few, would be established or highly constrained.

The Large Hadron Collider (LHC) would mainly act as a discovery machine like all other hadron colliders, and it is expected that a new  $e^+e^-$  Linear Collider (LC) [3] would complement the discoveries at LHC and make possible precision measurements of the parameter space governing the new physics scenario. We shall focus on a few of the many promising candidates of beyond standard model (BSM) physics scenarios. We look at their discovery prospects at future linear colliders through a particular channel of production, in association with photons in the final state. In section 2 we discuss the nice features of the associated photon signals at the next generation linear colliders and also try and motivate the idea to look for new physics signals through this production mode. In section 3 we take up different new physics models and discuss the proposed signal in isolating its signatures. We finally summarize and present our conclusions in section 4.

## 2 Associated Photon Signals at Linear Colliders

The history of associated photon signals to look for new physics signals goes way back [4–7] and has been a process of great interest for the physics programme at LEP [8,9]. It is also expected to be an important mode for new physics signals for future linear  $e^+e^-$  colliders [8]. The production of one or more photons in the final state along with the electroweak gauge bosons of the SM has been used extensively to probe the gauge interactions and the anomalous nature of photon interactions with other gauge bosons in the SM [10,11]. These studies have looked for new physics effects by looking at trilinear and quartic self-interactions of the gauge bosons involving photons. Precise knowledge of such interaction would help us understand the gauge

structure of the underlying physics. The most important result provided by single photon events is however the precise measurements of the number of light neutrino types [12, 13]  $N_\nu$  which is obtained by measuring the cross section of the process  $e^+e^- \rightarrow \gamma\nu\bar{\nu}$ . This process is also an important mode for new physics searches as it is sensitive to contributions from physics beyond the SM [14, 15].

We shall focus on the particular process

$$ee \rightarrow \gamma + X \quad (1)$$

where  $X$  can be any weakly interacting massive particle belonging to scenarios of new physics beyond the SM. We choose not to write the charges on the colliding particles of the beams, as we will consider two different cases, where the collision is either of electron-positron beams or electron-electron beams. The particle produced in association with  $\gamma$  carries two units of charge  $e$  when the collision is between two electron beams. The above process can be a very efficient tool to search for new physics signal as we show through the examples in the next section. The case of “photon+missing energy” is the most widely studied signal, which within the framework of SM accounts for the neutrino-counting experiment. It also gives an independent probe for the  $\gamma W^+W^-$  coupling at the high energy colliders as the contribution from the  $W$ -boson exchange becomes dominant. Thus it serves as an efficient tool to study anomalous couplings of the photon with the  $W$  boson as well. Several studies exist in the literature which consider the “photon+missing energy” process at future linear  $e^+e^-$  colliders to look for new physics signals [16–21].

It is needless to say that linear colliders will have the ability to make precise test of the structure of electroweak interactions at very short distances. Looking at the simplest process of  $e^+e^- \rightarrow f\bar{f}$ , the SM cross-section prediction can be put in the form

$$\frac{d\sigma}{d\cos\theta} \quad (e_L^- e_R^+ \rightarrow f_L \bar{f}_R) = \frac{\pi\alpha^2}{2s} N_C \cdot \left| Q_f + \frac{(\frac{1}{2} - \sin^2\theta_w)(I_f^3 - Q_f \sin^2\theta_w)}{\cos^2\theta_w \sin^2\theta_w} \frac{s}{s - m_Z^2} \right|^2 \cdot (1 + \cos\theta)^2 \quad (2)$$

where  $N_C = 1$  for leptons and 3 times for quarks,  $I_f^3$  is the weak isospin of  $f_L$ , and  $Q_f$  is the electric charge. For  $f_L$  production, the  $Z$  contribution typically interferes with the photon constructively for an  $e_L^-$  beam and destructively for an  $e_R^-$  beam. Thus, initial-state polarization is a useful diagnostic at the LC. Applied to familiar particles, they would provide a diagnostic of the electroweak exchanges that might reveal new heavy weak bosons or other types of new interactions. Simple annihilation processes can also be used to test for new interactions. However the best option to study such electroweak exchanges would be to study their physics at its resonance. The obvious reason being that, off-shell contributions will be strongly propagator-suppressed and suppress the new physics signal drastically. Our motivation for considering the process given in Eq. (1) is principally based on the fact that design for the future linear collider allow the machine to run at one or a few fixed center-of-mass energies. Single photon signals will allow the on-shell production of a massive particle  $X$  as long as  $M_X < \sqrt{s}$ . We show that this process can have a resonant production for  $X$  which shows up in the energy distribution of the photon. The idea is that as the photon carries away a variable amount of energy, it is possible for the remaining system (assuming  $X$  decays to some final states)

to strike a  $s$ -channel resonance of the particle  $X$ , just as initial state radiation (ISR) at LEP-2 has been seen to cause a ‘radiative return to the  $Z$ -boson pole. Then one can expect additional bump(s) over the continuum SM background in the photon energy distribution. The photon energy will be uniquely fixed by the well-known formula

$$E_\gamma = \frac{s - M_X^2}{2\sqrt{s}} \quad (3)$$

This signal is particularly interesting because of its simplicity and cleanliness. In the next section we discuss different physics models beyond SM and how they can leave their imprint on the “associated photon” signals at the future linear colliders.

### 3 New Physics and Associated Photon Signals

In this section we highlight the use of associated photon signals as a search tool for resonances at future linear colliders for massive particles predicted in various new physics models whose mass is less than the center-of-mass energy ( $\sqrt{s}$ ) of the machine. Such resonances are very likely to be missed if they are produced off-shell in the  $s$ -channel, and if their mass is also quite less than the ( $\sqrt{s}$ ) of the machine. We show that an associated photon carries the mass information of the produced particle in its energy distribution as given by Eq. (3). We investigate different models and study their features through the proposed signal.

#### 3.1 Supersymmetry

The most extensively studied new physics scenario over the last three decades has been supersymmetry (SUSY). If one considers  $R$ -parity ( $R_p$ ) conserving SUSY, ( $R$  is defined as  $R = (-)^{L+3B+2S}$ , where  $L, B$  and  $S$  stand, respectively, for the lepton number, baryon number and spin of a particle) then most of the search strategies are based on the fact that the lightest supersymmetric particle (LSP) is a massive weakly interacting neutral superparticle. It is stable and escapes detection and thus one expects large missing energy associated with SUSY signals. The analogue to our “single photon” signal in SUSY would be production of a pair of LSP’s with a photon in the final states which would invariably affect the cross section of the process  $e^+e^- \rightarrow \gamma \cancel{E}$  [16, 18, 19, 21]. However,  $R$ -parity can be violated [22] as long as either lepton number  $L$  or baryon number  $B$  if not both, is conserved. This can scramble the SUSY signals quite dramatically as the LSP is no longer stable and can decay within the detector. Admitting lepton number violating operators of the  $LL\bar{E}$  form (where we have assumed the conservation of baryon number  $B$ ), the relevant term in the superpotential can be written as

$$\mathcal{W}_{LL\bar{E}} = \lambda_{ijk} \epsilon_{ab} \hat{L}_i^a \hat{L}_j^b \hat{E}_k, \quad i, j = 1 \dots 3 \quad (4)$$

where  $\hat{L}_i \equiv (\hat{\nu}_{Li}, \hat{\ell}_{Li})^T$  and  $\hat{E}_i$  are the  $SU(2)$ -doublet and singlet superfields respectively whereas  $\epsilon_{ab}$  is the unit antisymmetric tensor. Clearly, the coupling constants  $\lambda_{ijk}$  are antisymmetric under the exchange of the first two indices; the 9 such independent couplings are usually labelled keeping  $i > j$ . Written in terms of

the component fields, the above superpotential leads to the interaction Lagrangian

$$\begin{aligned} \mathcal{L}_\lambda = \lambda_{ijk} & \left[ \tilde{\nu}^i \bar{\ell}_R^k \ell_L^j + \tilde{\ell}_L^j \bar{\ell}_R^k \nu_L^i + (\tilde{\ell}_R^k)^* \overline{(\nu_L^i)^c} \ell_L^j + (\tilde{\nu}^i)^* \bar{\ell}_L^j \ell_R^k + (\tilde{\ell}_L^j)^* \bar{\nu}_L^i \ell_L^k + \tilde{\ell}_R^k \bar{\ell}_L^j (\nu_L^i)^c \right. \\ & \left. - \tilde{\nu}^j \bar{\ell}_R^k \ell_L^i - \tilde{\ell}_L^i \bar{\ell}_R^k \nu_L^j - (\tilde{\ell}_R^k)^* \overline{(\nu_L^j)^c} \ell_L^i - (\tilde{\nu}^j)^* \bar{\ell}_L^i \ell_R^k - (\tilde{\ell}_L^i)^* \bar{\nu}_L^j \ell_L^k - \tilde{\ell}_R^k \bar{\ell}_L^i (\nu_L^j)^c \right]. \end{aligned} \quad (5)$$

The  $\mathcal{R}_p$  couplings are constrained by various experiments [23]. However for our case, we are interested in a process which deals with one particular type of coupling which allows single sneutrino production in association with a hard photon. The terms relevant for our discussion are the first and fourth ones on both first and second lines of Eq. (5), with  $j = k = 1$  on the first line and  $i = k = 1$  on the second.

The sneutrino decay width, which never rises above 3–4 GeV, allows us to apply the *narrow-width approximation* and therefore, we solely consider on-shell production of sneutrinos (of muonic or tauonic flavour). If, indeed,  $m_{\tilde{\nu}} < \sqrt{s}$ , then the cross-section for  $e^+ + e^- \longrightarrow \tilde{\nu}_{\mu/\tau}(\tilde{\nu}_{\mu/\tau}^*)$  will be strongly propagator-suppressed. As discussed in section 2 the processes of interest for us then becomes

$$\begin{aligned} e^+ + e^- & \longrightarrow \gamma + \tilde{\nu}_{\mu/\tau}(\tilde{\nu}_{\mu/\tau}^*) \longrightarrow \gamma + e^+ + e^- \\ & \hookrightarrow \gamma + \tilde{\nu}_{\mu/\tau}(\tilde{\nu}_{\mu/\tau}^*) \longrightarrow \gamma + \nu_{\mu/\tau}(\bar{\nu}_{\mu/\tau}) + \tilde{\chi}_{1/2/3/4}^0 \\ & \hookrightarrow \gamma + \tilde{\nu}_{\mu/\tau}(\tilde{\nu}_{\mu/\tau}^*) \longrightarrow \gamma + \mu^\mp/\tau^\mp + \tilde{\chi}_{1/2}^\pm, \end{aligned} \quad (6)$$

where the application of the narrow-width approximation ensures an almost monochromatic photon of energy given by Eq. (3), where  $X$  is to be replaced by the sneutrino mass. This, potentially, would stand out against the continuum spectrum arising from the SM background. Since the sneutrino  $\tilde{\nu}_{\mu/\tau}$  can have a variety of decay channels, we can simply tag on a hard isolated photon associated with any of these decay channels and look for a line spectrum superposed on the continuum background. This will lead to clear signals of sneutrino production. Moreover, the  $R$ -parity-violating decays of the sneutrino will set up multi-lepton final states (with associated photons) which will have little or no SM backgrounds worth considering. For such states a mono-energetic photon will clinch the issue of sneutrino production [24]. The specific reaction on which we focus in this case is the associated photon process

$$e^+ + e^- \longrightarrow \gamma + \tilde{\nu}_{\mu/\tau}(\tilde{\nu}_{\mu/\tau}^*) .$$

The squared and spin-averaged matrix element for this is, then

$$|\overline{\mathcal{M}}|^2 = 8\pi\alpha \lambda_{1j1}^2 \frac{s^2 + \tilde{m}_j^4}{tu} \theta(s - \tilde{m}_j^2) \quad (7)$$

where  $\tilde{m}_j$  is the mass of the muonic ( $j = 2$ ) or tauonic ( $j = 3$ ) sneutrino. The collinear singularity in Eq. (7), so characteristic of massless electrons and photons, is automatically taken care of once one imposes restrictions on the phase space commensurate with the detector acceptances. In the rest of the analysis, we shall require the photon to be sufficiently hard and transverse, namely

$$\begin{aligned} \text{pseudorapidity :} & \quad |\eta_\gamma| < \eta_\gamma^{(max)} = 2.0 , \\ \text{transverse momentum :} & \quad p_{T\gamma} > p_{T\gamma}^{(min)} = 20 \text{ GeV} . \end{aligned} \quad (8)$$

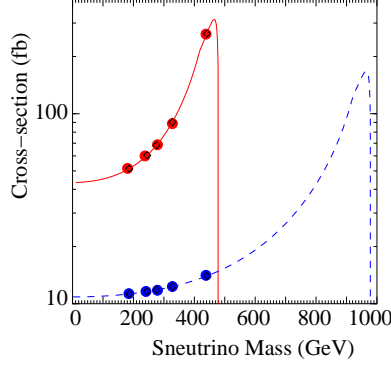


Figure 1: Cross sections for sneutrino production with associated photons at a linear collider for  $\lambda_{1j1} = 0.03$ . Solid red (dashed blue) lines correspond to a 500 GeV (1 TeV) center-of-mass energy. The cuts of Eq.(8) have been imposed. The points marked with bullets are for a  $\tilde{\nu}_\mu$  resonance at the Snowmass MSugra points 1a, 1b, 3, 4, and 5. At the point 2, the sneutrino is beyond the kinematic reach of the linear collider. If sneutrinos are not distinguished from anti-sneutrinos, the cross-section(s) would be doubled.

The cross-section is plotted in Fig. (1) as a function of  $\tilde{m}_j$  and with  $\lambda_{1j1} = 0.03$  for a linear collider running at (a) 500 GeV and (b) 1 TeV.

As the graph shows, we obtain cross-sections typically in the range 50 fb–250 fb. At a linear collider with around  $500 \text{ fb}^{-1}$  of integrated luminosity, this amounts to the production of a very large number of sneutrinos along with an associated monochromatic photon. Thus, even if  $\lambda_{1j1}$  were to be smaller by an order of magnitude, we would still have a fairly large number of such distinctive events. It is clear, therefore, that if the sneutrino is kinematically accessible to a linear collider, low statistics will not be the major hurdle in their detection.

We focus on the MSugra spectrum and, specifically, on the six representative points chosen at the 2001 Snowmass conference. The latter along with the spectrum are described (for  $\mu > 0$ ) in the table below.

Depending on the various decay modes available for the sneutrino, we focus on four classes of final states, which are (1)  $\gamma ee$ , (2)  $\gamma ee + \cancel{E}$ , (3)  $\gamma \ell_i \ell_j + \cancel{E}$  and (4)  $\gamma 4\ell + \cancel{E}$ . The first kind arises from the direct  $R$ -parity-violating decay of the sneutrino and would have a large SM background from radiative Bhabha scattering. The second and third ones are obviously reproduced by  $WW$ -production. The last type arises from higher-order effects in the SM and has very little background. This final state arises from the direct  $R$ -parity violating decay of the sneutrino into an  $e^+e^-$  pair, with, of course an associated photon from the initial state. The branching ratio of the sneutrino to this mode is quite significant for  $\lambda_{1j1} \sim 0.03$  and hence the signal has a reasonable cross-section. We carry out our analysis at a 500 GeV  $e^+e^-$  collider. To detect this final state, we impose a set of acceptance sets, namely that each of the particles must not be too close to the beam pipe,

$$|\eta(e^\pm)|, |\eta(\gamma)| < 2.0 \quad (9)$$

and that they should carry sufficient transverse momenta

$$p_T(e^\pm) > 10 \text{ GeV} \quad \text{and} \quad p_T(\gamma) > 20 \text{ GeV} . \quad (10)$$

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|             |                           |                              |                           |                   |
|-------------|---------------------------|------------------------------|---------------------------|-------------------|
| <b>1a :</b> | $M_0 = 100 \text{ GeV},$  | $m_{1/2} = 250 \text{ GeV},$ | $A_0 = -100 \text{ GeV},$ | $\tan \beta = 10$ |
| <b>1b :</b> | $M_0 = 200 \text{ GeV},$  | $m_{1/2} = 400 \text{ GeV},$ | $A_0 = 0,$                | $\tan \beta = 30$ |
| <b>2 :</b>  | $M_0 = 1450 \text{ GeV},$ | $m_{1/2} = 300 \text{ GeV},$ | $A_0 = 0,$                | $\tan \beta = 10$ |
| <b>3 :</b>  | $M_0 = 90 \text{ GeV},$   | $m_{1/2} = 400 \text{ GeV},$ | $A_0 = 0,$                | $\tan \beta = 10$ |
| <b>4 :</b>  | $M_0 = 400 \text{ GeV},$  | $m_{1/2} = 300 \text{ GeV},$ | $A_0 = 0,$                | $\tan \beta = 50$ |
| <b>5 :</b>  | $M_0 = 150 \text{ GeV},$  | $m_{1/2} = 300 \text{ GeV},$ | $A_0 = -1 \text{ TeV},$   | $\tan \beta = 5$  |

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| Point | $\tilde{\nu}_\mu$ | $\tilde{\nu}_\tau$ | $\tilde{\tau}_1$ | $\tilde{\tau}_2$ | $\tilde{\chi}_1^0$ | $\tilde{\chi}_2^0$ | $\tilde{\chi}_1^\pm$ |
|-------|-------------------|--------------------|------------------|------------------|--------------------|--------------------|----------------------|
| 1a    | 186               | 185                | 133              | 206              | 96                 | 177                | 176                  |
| 1b    | 328               | 317                | 196              | 344              | 160                | 299                | 299                  |
| 2     | 1454              | 1448               | 1439             | 1450             | 80                 | 135                | 104                  |
| 3     | 276               | 275                | 171              | 289              | 161                | 297                | 297                  |
| 4     | 441               | 389                | 268              | 415              | 119                | 218                | 218                  |
| 5     | 245               | 242                | 181              | 258              | 120                | 226                | 226                  |

In addition, each pair of the final state particles should be well separated:

$$\delta R > 0.2, \quad (11)$$

where  $(\delta R)^2 \equiv (\Delta\phi)^2 + (\Delta\eta)^2$  with  $\Delta\eta$  and  $\Delta\phi$  respectively denoting the separation in rapidity and azimuthal angle. On analyzing the distributions in various kinematic variables, it is found that an additional rapidity cut on the difference of the rapidity variable of the final state electrons ( $|\Delta\eta_{ee}| = |\eta_{e-} - \eta_{e+}|$ ) suppresses the huge SM background which comes from the t-channel dominated Bhabha scattering without reducing the signal by much. A detailed analysis helps us to fix this value at 1.7 which we impose for our analysis. The remaining classes of final states do not have large SM background and are not affected by the cuts too much. So for the leptons and the photon, we choose the cuts to be the same as before, namely those listed in Eqs.(9–11). In addition, we demand that the missing transverse momentum be sufficiently large, *viz.*

$$\cancel{p}_T > 20 \text{ GeV} \quad (12)$$

for it to be considered a genuine physics effect.

We now focus on the photon in the final state which is our main trigger. We show the distribution in photon energy for the first two classes of final states in Fig. (2) The corresponding fluctuation (Gaussian) in the SM is also shown at 1, 3 and 5 standard deviations. In Fig. (2a), the clear peaks in the energy distribution of the photon, over the continuum SM background gives clear hints of the sneutrino production. The mass of the sneutrino can also be very easily determined by the formula given in Eq. (3). However, for class (1), we also have the leptons to trigger at and whose invariant mass will reconstruct the mass of the sneutrino. But the efficacy of the photon signal becomes clear when we look at more complicated decays of the sneutrino. In this case the final state leptons cannot reconstruct the sneutrino mass, because of the

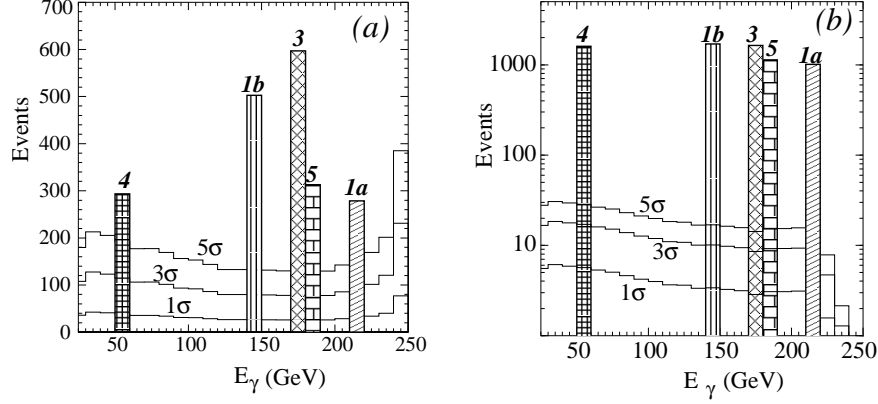


Figure 2: The photon energy distribution when the final states are (a)  $\gamma e^+ e^-$  and (b)  $\gamma e^+ e^- \cancel{E}$  at a 500 GeV  $e^+ e^-$  collider. The integrated luminosity is  $\mathcal{L} = 100 \text{ fb}^{-1}$ . The different representative Snowmass points are shaded in different patterns.

presence of neutrinos in the final state. But the peaks in the photon energy distribution will be a complete give-away for the production of sneutrino. This feature gets highlighted in Fig. (2b) where the final state has missing energy associated with it. The same feature will repeat itself for the other final states. Thus we find that even with complicated decay modes for the sneutinos, which render the reconstruction of the parent particles mass quite improbable, the clear peaks in the photon energy distribution will save the situation.

### 3.2 Extra Dimensions

Theories with extra dimensions as possible solution to the hierarchy problem and unification have recently attracted enormous attention and interest. We focus on two class of models *viz.* ADD [25] and RS [26] model. In the ADD model the Standard Model fields are confined on a (1+3)-dimensional subspace ( $D_3$  brane) of a  $(1 + 3 + d)$ -dimensional spacetime (bulk). The  $d$  extra dimensions are compactified, typically on a  $d$ -torus of radius  $R_c$ . Gravity which reflects the geometry of spacetime itself cannot be confined and is free to move in the bulk. The essential idea contained in their model was that when we match the higher dimensional theory with the 4D effective theory, the following relation is obtained

$$M_{Pl}^2 = M_S^{d+2} V_{(d)} = M_S^{d+2} (2\pi R_c)^d \quad (13)$$

where  $M_S$  is the fundamental Planck scale (in the bulk),  $M_{Pl}$  is the observed 4D Planck scale  $\sim 10^{19}$  GeV and  $V_{(d)}$  is the volume of the extradimensional space. Now if  $R_c > 1/M_{Pl}$ , then the fundamental Planck scale  $M_*$  will be lowered from  $M_{Pl}$ . Under this assumption, the fundamental Planck scale can be brought down to values as low as  $M_S \sim 1$  TeV, consistent with present experimental bounds. Thus it is able to resolve the hierarchy problem by cutting of the SM at the TeV scale. The Kaluza Klein (KK) excitations of the (bulk) graviton are very closed spaced with each mode separated by  $1/R_c$  which couple to the SM particles by a coupling strength  $\sim 1/M_{Pl}$ . However there are huge number of these excitations which



contribute to collectively build up observable effects at the electroweak scale. Thus one hopes to see their effect at current and future experiments. Another approach is shown in the RS [26] model, which resolves the hierarchy problem, even with small extra dimension. They in fact choose a different spacetime geometry (non-factorisable metric) and an additional 3-brane in their argument with only one extra dimension added to our (1+3)-dimensional spacetime. In contrast to the ADD relation Eq. (13), the 4-dimensional Planck scale in the RS approach is:

$$M_{Pl}^2 = \frac{M_5^3}{k} [1 - e^{-2kR_c\pi}] \quad (14)$$

where  $M_5$  is the fundamental scale of the model,  $R_c$  the compactification size and  $k$  determines the curvature of the space. A TeV energy scale can be generated from the 4-dimensional Planck scale if  $kR_c \sim 12$ , and thus providing a solution to the hierarchy problem between the electroweak scale of the standard model and the 4-dimensional Planck scale.

The important differences of the RS model with the ADD model are

- Each KK excitation of the bulk graviton has a mass

$$M_n = x_n \mathcal{K} e^{-\mathcal{K}R_c\pi} \equiv x_n m_0 \quad (15)$$

where  $m_0 = \mathcal{K} e^{-\mathcal{K}R_c\pi} \sim 100$  GeV is the graviton mass scale and  $x_n$  are the zeros of the Bessel function  $J_1(x)$  of order unity ( $n \in \mathbf{Z}$ ). This means that the Kaluza-Klein gravitons have masses of a few hundred GeV, unlike the ADD case, where the masses start from  $\sim 1$   $\mu$ eV.

- Each Kaluza-Klein excitation of the bulk graviton couples to matter as [27–29]

$$\kappa e^{\mathcal{K}R_c\pi} = \frac{4\sqrt{\pi}}{M_P} e^{\mathcal{K}R_c\pi} \equiv \frac{4\sqrt{\pi}c_0}{m_0} \quad (16)$$

where  $\kappa = \sqrt{16\pi G_N}$  and  $c_0 = \mathcal{K}/M_P \simeq 0.01 - 0.1$  is an effective coupling constant, whose magnitude is fixed by (a) naturalness and (b) requiring the curvature of the fifth dimension to be small enough to consider linearized gravity on the ‘visible’ brane.

RS gravitons, thus, resemble weakly-interacting massive particles (WIMPs) in most models, except for (a) the fact that there always exists a tower of graviton Kaluza-Klein modes and (b) these are spin-2 particles. In phenomenological studies of the RS model, the mass scale  $m_0$  and the ratio  $c_0$  may be treated as free parameters<sup>1</sup>: they are convenient replacements for the fundamental quantities  $\mathcal{K}$  and  $R_c$ :

$$\mathcal{K} = c_0 M_P, \quad R_c = \frac{1}{\pi \mathcal{K}} \log \frac{\mathcal{K}}{m_0} \quad (17)$$

We show that “single photon” signals can be very effectively used to look for such (WIMP) using the radiative return technique. The motivation remains similar to the exercise done for the sneutrino search. However, a major difference is that KK gravitons in the RS model follow a distinct relation according to the zeros of the Bessel function  $J_1(x)$  of order unity ( $n \in \mathbf{Z}$ ). Thus, by exciting multiple resonances in the photon energy

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<sup>1</sup>The alternative choice of  $\Lambda_\pi = \overline{M}_P e^{-\mathcal{K}R_c\pi} = m_0/\sqrt{8\pi}c_0$  instead of  $m_0$  and of  $\mathcal{K}/\overline{M}_P = \sqrt{8\pi}c_0$  instead of  $c_0$  may also be found in the literature [29].

distribution, one could in fact check this relation from the appearance of position of the peaks. This gives a very strong evidence for RS-type scenario.

Since the RS gravitons have widths which grow with increasing value of  $c_0$ , they might not be narrow. It can be written down as

$$\Gamma_n = c_0^2 x_n^3 m_0 \sum_P \Delta_{P\bar{P}}^{(n)} \quad (18)$$

where the sum  $\sum_P$  runs over all pairs of particles ( $G_n \rightarrow P\bar{P}$ ) and the  $\Delta_{P\bar{P}}^{(n)}$  are dimensionless functions of  $x_n$  and the ratios  $r_P = m_P/m_0$ . For details one can look at Ref. [20]. We calculate the full  $2 \rightarrow 3$  process for  $e^+e^- \rightarrow \gamma\nu\bar{\nu}$ . So here our final state is just a single hard photon with unbalanced (missing) energy. Studies with other mode of decay of the graviton can be found in Ref [30]. We choose the invisible decay mode of the RS gravitons, as this would also enable us to compare the signal with the ADD signal, where the equivalent process is  $e^+e^- \rightarrow \gamma G_n$ , and the ADD graviton escapes detection and carries the missing energy, because of its very small coupling to SM particles.

We present results for two values of center-of-mass energy, viz.,  $\sqrt{s} = 1$  TeV and  $\sqrt{s} = 2$  TeV. Noting that the final state consists of a single hard isolated photon, we impose the following kinematic cuts

- The photon should have energy  $E_\gamma \geq 20$  GeV.
- The photon scattering angle  $\theta_\gamma$  should satisfy  $15^\circ \leq \theta_\gamma \leq 165^\circ$ .

These ensure that the tagged photon does not arise from beamstrahlung or other similar sources. Since the dominant SM background for the “single photon” signal comes from the exchange of  $W$ -bosons in the  $t$ -channel, use of polarized beams (right-polarized electron beam and left-polarized positron beam), suppresses the background very efficiently and enhances the signal, when compared to the case of unpolarized beams. We use the partial polarization of the initial beams to be  $\mathcal{P}_e = 0.8$  and  $\mathcal{P}_p = -0.6$  for the presented analysis. In Fig. (3) we show the energy distribution of the single hard photon in the final state. The dashed (black) line represents the background (note the  $Z$ -boson peak at the extreme right), while the remaining curves correspond to the signal for low ( $-\cdot-$ ) and high (solid) values of the parameter  $c_0$ . Sharp resonances are obtained with the parameter choice  $c_0 = 0.01$  and  $m_0 = 125$  GeV. It corresponds to graviton resonances with  $M_n \simeq 479, 877, 1272$  and  $1665$  GeV for  $n = 1, 2, 3$  and  $4$  respectively. Only the first two are kinematically accessible at a 1 TeV machine, but all four will be accessible if the center-of-mass energy rises to 2 TeV. Observe that the resonance peaks broaden as the order  $n = 1, 2, \dots$  of the Kaluza-Klein excitation increases<sup>2</sup>. In the upper halves of the two graphs in Fig. (3), we display the differential cross-section for the process  $e^+e^- \rightarrow \gamma + \cancel{E}$ . The bottom halves show the same distribution, except that now we exhibit the signal-to-background (S/B) ratio. Not only does this remove the uninteresting  $Z$ -peak, but it also takes care of any radiative corrections, efficiency factors, etc, which can be written in a factorisable form.

Two (or more) clear resonances seen in the photon spectrum, or rather, in the signal-to-background ratio would correspond to a relatively low value of  $m_0$  and a small value of  $c_0$ , and would constitute a strong hint

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<sup>2</sup>The present bounds from Tevatron data [31] rule out the low lying KK modes ( $M_n > 800$  GeV) for the multiple resonances to be seen at a 1 TeV machine. However the characteristic features can be still seen at the 2 TeV machine or CLIC with 3 TeV center-of-mass energy

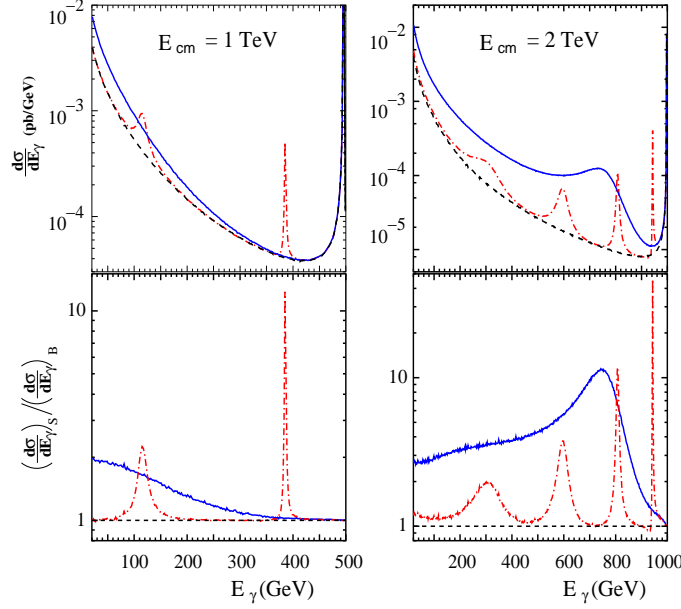


Figure 3: Energy spectrum of the tagged photon for the choices  $m_0 = 125$  GeV,  $c_0 = 0.01$  in red ( $-\cdot-$ ) corresponding to narrow resonance(s) and  $m_0 = 250$  GeV,  $c_0 = 0.07$  in blue (solid) corresponding to broad, indistinct resonances. In the ordinate labels, S and B denote signal (SM plus gravitons) and background (SM only) respectively. Note that the right-most peak (almost flush with the edge of the box) in the upper graphs, which is due to the Z-boson, can be removed by taking the S/B ratio.

of RS gravity. For a strong confirmation, positions of the two resonances will bear a definite relation if they are due to RS gravitons. Using Equation (3) this works out to the requirement that, for resonance values  $E_\gamma^{(1)} > E_\gamma^{(2)}$ ,

$$\sqrt{\frac{\sqrt{s} - 2E_\gamma^{(1)}}{\sqrt{s} - 2E_\gamma^{(2)}}} = \frac{M_1}{M_2} = \frac{x_1}{x_2} \simeq 0.546. \quad (19)$$

It would be a very remarkable coincidence, indeed, if some other form of new physics reproduces such a relation.

A single sharp resonance will allow a consistent fit to the two-parameter RS model. One can then look at other distributions, in addition to the single photon events to decipher the underlying physics. Just to compare with the case of SUSY, a single resonance rules out the conventional  $R$ -parity conserving SUSY. Even with  $R$ -parity violation, one can use the angular distribution of the leptons to differentiate a spin-0 resonance from that of spin-2 graviton resonance.

When the mass  $m_0$  and/or the coupling  $c_0$  is large, and no clear resonance structure is discernible, is it possible to distinguish between the  $\gamma + \cancel{E}$  signal arising from RS graviton from those arising from ADD gravitons [27, 32]? The energy distribution of the photon will just show an excess without any bumps (which was the give away distinction from the ADD smooth spectrum) in the continuous photon energy spectra. However, if one consider this process *in conjunction* with a benchmark process, like  $e^+e^- \rightarrow \mu^+\mu^-$ , for example, we do find a marked difference. A correlation plot showing the cross-section for the single photon

signal vis-à-vis the cross-section for muon pair-production proves to be clinical in distinguishing the two different models [20].

### 3.3 Other exotics

In this section we consider our “associated photon” signal at a linear  $e^-e^-$  collider to look for doubly charged Higgs bosons which arise in a number of physics scenarios [33–39], the most common models to accommodate such scalars are those with triplet Higgs. An added feature often associated with doubly-charged Higgs is the possibility of  $\Delta L = 2$  couplings with leptons which can be very effectively probed at  $e^-e^-$  colliders and give a strong motivation towards running the future linear collider in this mode. The  $\Delta L = 2$  coupling appears in the Lagrangian as

$$L_Y = ih_{ij}\Psi_{iL}^T C\tau_2\Phi\Psi_{jL} + h.c. \quad (20)$$

where  $i, j = e, \mu, \tau$  are generation indices, the  $\Psi$ ’s are the two-component left-handed lepton fields, and  $\Phi$  is the triplet with  $Y = 2$  weak hypercharge. This leads to mass terms for neutrinos once the neutral component  $\phi^0$  of  $\Phi$  acquires a vacuum expectation value (vev). Constraints on the  $\rho$ -parameter puts strong limits on the triplet vev translating into limits on the L-violation Yukawa couplings which constrain the collider signals for doubly-charged scalars sought through  $\Delta L = 2$  interactions.

At linear  $e^+e^-$  colliders, such exotics would have to be produced in pair, which limits the available phase space for its production. However, they can be produced singly if the collision is between electron beams. As before, our ignorance of the doubly charged Higgs mass does not allow its production at resonance. We point out the usefulness of looking for doubly-charged scalars in an  $e^-e^-$  collider, in the radiative production channel. Not only does this allow us to use the “radiative return” technique, but also extends the phase space for its on-shell production, constrained only by kinematic cuts on the photon and the machine energy. With this in view, we consider the process [40–42]

$$e^-e^- \longrightarrow \phi^{--}\gamma \longrightarrow Y\gamma$$

at a  $\sqrt{s} = 1$  TeV  $e^-e^-$  machine, concentrating on the hard single photon in the final state. Here  $Y$  represents the decay products of the doubly charged scalar. The hard photon in the final state will be monochromatic if a doubly-charged resonance is produced, irrespective of what it decays into. For our analysis, taking the radiative production of the scalar  $\phi^{--}$  as the benchmark process, we concentrate only on the flavor diagonal coupling  $h_{ee}$ , which we choose to be  $h_{ee} = 0.1$  which respects the most stringent bounds coming from muonium-antimuonium conversion results which for flavor diagonal coupling is  $h < 0.44 M_\phi^{\pm\pm} \text{ TeV}^{-1}$  at 90% C.L [43]. As shown before, the on-shell radiative production of a doubly-charged scalar gives an almost monochromatic photon of energy given by Eq. (1).

The major SM background that contributes to the above process is the radiative Moller scattering process:  $e^- + e^- \rightarrow \gamma + e^- + e^-$  which, although a continuum background, is quite large. The event selection criteria largely aims at suppressing this continuum background. We impose the following set of cuts.

- Rapidity cut on the final state particles:  $|\eta(e^-)| < 1.5$  and  $|\eta(\gamma)| < 2.5$

- minimum cut on energies:  $E(\gamma) > 20$  GeV,  $E(e^-) > 5$  GeV
- to ensure that the final state particles are well separated in space for the detectors to resolve events :  $\delta R > 0.2$

Using the above cuts we make an estimate of the SM background and the signal. We focus on the main trigger, *viz.* the photon. In Fig 4(a) we show the distribution of the photon energy, where we have superposed the differential cross-section for *signal+background* in each bin over the SM background. A pronounced peak can be seen in the photon energy distribution, due to the monochromaticity of the photon, corresponding to the recoil energy against the scalar resonance through the relation of Eq. (1). To make our analysis realistic, we have smeared the photon energy by a Gaussian function whose half-width is guided by the resolution of the electromagnetic calorimeter [3] and also incorporated the effects of ISR which often results in substantial broadening of the peak. We show the resulting peak for three choices of scalar mass (300, 600, 900 GeV). In Fig 4(a) we only look at the final state hard transverse photon in  $e^-e^- \rightarrow \gamma + \phi^{--} \rightarrow \gamma + Y(\text{anything})$ .

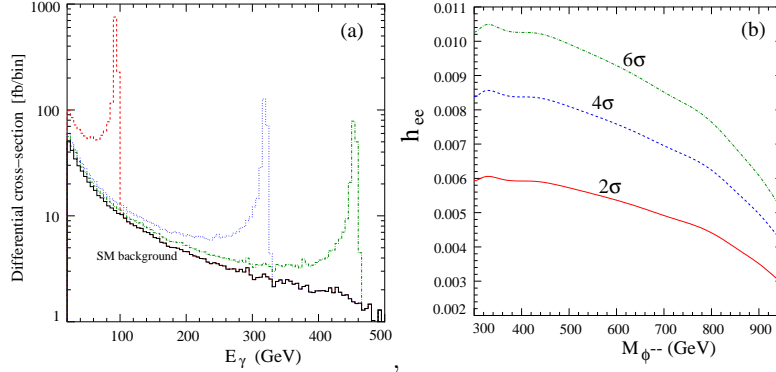


Figure 4: **(a)** Differential cross-sections against photon energy  $E_\gamma$  when  $\phi^{--} \rightarrow Y(\text{anything})$ . The dash-dot-dash (green) line corresponds to  $M_{\phi^{--}} = 300$  GeV, dotted (blue) line corresponds to  $M_{\phi^{--}} = 600$  GeV and the dashed (red) line corresponds to  $M_{\phi^{--}} = 900$  GeV respectively. **(b)** Illustrating the reach of the coupling constant at which the resonances in the  $E_\gamma$  distribution can be identified over the fluctuations in the SM background. The assumed luminosity is  $100 \text{ fb}^{-1}$ .

The distribution again shows peaks corresponding to the recoil against the massive scalars, irrespective of the knowledge of the decay products of the scalar. In fact the signal here receives a relative boost as it is not suppressed by considering any further decay since the  $BR(\phi^{--} \rightarrow Y) = 100\%$ . The fact that looking at a single photon against the backdrop of a continuum background makes it possible to identify a LFV ( $\Delta L = 2$ ) process in a model independent way, makes this signal worth studying at a future  $e^-e^-$  collider and running the linear collider in this mode. One can also constrain the strength of the  $eeH$  coupling. Since the rates for the signal depend directly on the  $eeH$  coupling squared, in Fig 4(b) we show the strength of the coupling for which the peaks would stand out against the fluctuations in the SM background. In our analysis we have assumed a luminosity of  $\mathcal{L} = 100 \text{ fb}^{-1}$ . The fact that we are not looking at any specific final state arising from  $\phi^{--}$  decay improves the reach of this search channel.

## 4 Summary

To summarize, the cleanliness of central photon detection at a high energy linear collider can be very helpful in identifying a weakly interacting massive particles predicted in various new physics scenarios beyond the SM. To highlight this, we have considered three different scenarios of physics beyond the SM and shown how a hard photon in the final state can be useful in extracting the underlying physics by looking at the photon energy distribution. The peaks in the hard photon energy can be helpful in two ways. First, one does not need to tune the two electron beams, and can therefore work without a prior knowledge of the particles ( $X$ ) mass against which the photon recoils. Secondly, this method is shown to work even if the  $X$  dominantly decays into states that are not clean enough for the resonance to be identified. Thus, as soon as one succeeds in reducing the SM backgrounds, one can clearly see non-SM interactions, just by looking at the accompanying hard photon. More exotic resonances such as extra  $Z'$  bosons predicted in models beyond SM with extra  $U(1)$ 's, KK gauge bosons predicted in the universal extra dimension models (UED), bileptons, etc., are amenable to detection in this manner. Although we have focussed primarily on resonant searches, the associated photon signals are also sensitive to new physics signals [14–17] in other forms.

Another way of using associated photons to capture resonances at future linear colliders would be through beamstrahlung and ISR photons, by looking at the simple scattering process  $ee \rightarrow X^* \rightarrow PP$  where  $P$  is any SM particle. The radiated photons will cause a spread in the beam energy and cause some of the events to take place at values lower than the actual center-of-mass energies ( $\sqrt{s}$ ). If the exchange particle mass is less than  $\sqrt{s}$  and the events due to radiation happen around the resonance(s), it will provide a huge enhancement in the cross-section [44, 45]. Thus, associated photon signals will play an important role at the next generation linear colliders and the radiative return technique proves to be a crucial tool, both in the case of tagged (hard) photons or radiative photons which go down the beam pipe, but may cause large energy spread to excite new resonances in the invariant mass distributions of the final state. The physics of “associated/single photon” signals at the next generation of linear colliders merits a lot of attention and can prove to be an important tool to study physics beyond the SM.

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